

Top-quark rare decay $t \rightarrow ch$ in R-parity-violating SUSYG. Eilam¹, A. Gemintern¹, T. Han², J. M. Yang³ and X. Zhang⁴¹ *Physics Department, Technion, 32000 Haifa, Israel*² *Department of Physics, University of Wisconsin, Madison, WI 53706, USA*³ *Institute of Theoretical Physics, Academia Sinica, Beijing 100080, China*⁴ *Institute of High Energy Physics, Academia Sinica, Beijing 100039, China***Abstract**

The flavor-changing top-quark decay $t \rightarrow ch$, where h is the lightest CP-even Higgs boson in the minimal supersymmetric standard model, is examined in the R-parity-violating supersymmetric model. Within the existing bounds on the relevant R-parity-violating couplings, the branching fraction for $t \rightarrow ch$ can be as large as about 10^{-5} in some part of the parameter space.

The study of heavy-particle decays via flavor-changing neutral-currents (FCNC) has been playing an important role in testing the standard model (SM) and probing new physics beyond the SM. As the heaviest elementary particle in the SM with a mass at the electroweak scale, the top quark is more likely to be sensitive to new physics. Kinematically it is accessible to many FCNC decay modes, such as $t \rightarrow cV$ ($V = \gamma, Z, g$) and $t \rightarrow ch$, where h is a Higgs boson. In the SM these FCNC decay modes are highly suppressed by GIM mechanism, with branching fractions typically of $10^{-13} - 10^{-10}$ [1,2], which are too small to be detectable at collider experiments. On the other hand, observation of any of such FCNC top-quark decays would be robust evidence for new physics [3–6].

Top quarks will be copiously produced at the next generation of hadron colliders. At the upgraded Fermilab Tevatron with an integrated luminosity of 10 fb^{-1} , there will be about 8×10^4 top quarks produced, while there will be about 100 times more at the LHC with the same luminosity. With such large data samples, good sensitivities may be reached for searching for the rare decay channels $t \rightarrow cV$ [3], and for studying other related processes at hadron colliders [4]. A more recent study showed that the channel $t \rightarrow ch$ could also be detectable [5], reaching a sensitivity level for the branching fraction $Br(t \rightarrow ch) \sim 5 \times 10^{-5}$ at the LHC and a few percent at the Tevatron. While these high detection sensitivities are still far above the SM expectation for the rare decay channels [1,2], in many scenarios beyond the SM the branching fractions of these FCNC top-quark decays could be significantly enhanced [1,6–11].

In the minimal supersymmetric (SUSY) standard model (MSSM) [12] with R-parity conservation, it was shown [8] that the possibility for observing the decay channel $t \rightarrow ch$ could be greatly enhanced (here h is the lightest CP-even Higgs boson). Kinematically this decay mode is always allowed because of the strict theoretical upper bound on the Higgs boson mass [13], and the decay receives dominant contributions from the SUSY QCD loops of flavor-changing interactions [8]. If the gluino and squarks involved in the contributing SUSY QCD loops are both light of order 100 GeV, the branching fraction could be enhanced to a level of 10^{-5} . The branching fraction falls off quickly for heavier sparticles in the loops.

In this Letter we examine the R-parity-violating contributions to $t \rightarrow ch$. It is well-known that the R-parity conservation in SUSY theory, which implies the separate conservation of baryon number and lepton number, is put in by hand. R-parity violation (\mathcal{R}) can be made perfectly consistent with other fundamental principles such as gauge invariance, supersymmetry and renormalizability [14]. If R-parity violation is included in the MSSM, $t \rightarrow ch$ will receive new contributions from the loops of R-parity-violating interactions. Such contributions can be significant for the following reasons. First, such contributing loops only involve a single sparticle, *i. e.*, a squark or a slepton. The mass suppression in the loops thus becomes less severe than that in the MSSM. Second, the relevant \mathcal{R} couplings inducing $t \rightarrow ch$ involve the third-generation fermions and are subject to rather weak bounds from low-energy experiments. We find that the branching fraction for $t \rightarrow ch$ from \mathcal{R} contributions can be indeed as high as about 10^{-5} in some part of the parameter space, reaching a level of potentially accessible at the LHC with a high luminosity.

We start our study by writing down the \mathcal{R} superpotential of the MSSM

$$\mathcal{W}_{\mathcal{R}} = \frac{1}{2} \lambda_{ijk} L_i L_j E_k^c + \lambda'_{ijk} L_i Q_j D_k^c + \frac{1}{2} \lambda''_{ijk} U_i^c D_j^c D_k^c + \mu_i L_i H_2, \quad (1)$$

where $L_i(Q_i)$ and $E_i(U_i, D_i)$ are the left-handed lepton (quark) doublet and right-handed

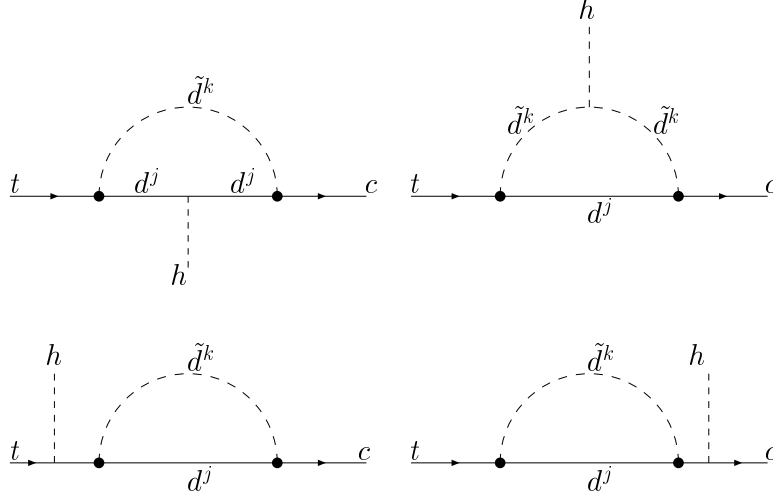


FIG. 1. Feynman diagrams of $t \rightarrow ch$ induced by \mathcal{B} SUSY interactions.

lepton (quark) singlet chiral superfields. i, j, k are generation indices and c denotes charge conjugation. Note that $SU(2)_L$ and $SU(3)_C$ indices have been suppressed. $H_{1,2}$ are the Higgs-doublets chiral superfields. The λ_{ijk} and λ'_{ijk} are the lepton-number-violating (\mathcal{L}) couplings, and λ''_{ijk} the baryon-number-violating (\mathcal{B}) couplings. Constraints on these couplings have been obtained from various low-energy processes [15–21] and their phenomenology at hadron and lepton colliders have been intensively investigated recently [21,22]. Note that although it is theoretically possible to have both \mathcal{B} and \mathcal{L} interactions, the non-observation of proton decay prohibits their simultaneous presence, at least in the first two fermion generations. We therefore assume the existence of either \mathcal{L} couplings or \mathcal{B} couplings, and investigate their separate effects in top-quark decay $t \rightarrow ch$.

We focus our attention only on the tri-linear supersymmetric \mathcal{R} interactions in Eq. (1) and assume that the bi-linear terms $\mu_i L_i H_2$ can be rotated away by a field redefinition [14]. In this case the FCNC decay $t \rightarrow ch$ is induced by only the tri-linear \mathcal{R} interactions via loops. Note that, in principle, there are also possible \mathcal{R} terms in the soft-breaking part [23]. In that case, it is no longer possible in general to rotate away the bi-linear terms [23,24] and such bi-linear terms will cause the mixing between the neutral Higgs bosons and the sneutrinos ($\tilde{\nu}$). As studied in [10], the FCNC decay $t \rightarrow c\tilde{\nu}$ can be induced by the one-loop diagrams of the tri-linear couplings. Followed by the oscillation $\tilde{\nu} \rightarrow h$ induced by the bi-linear \mathcal{R} terms, one can also have $t \rightarrow c\tilde{\nu} \rightarrow ch$, which, however, is only appreciable when h and $\tilde{\nu}$ are nearly degenerate. We will not consider this possibility further.

In terms of the four-component Dirac notation, the Lagrangian of the \mathcal{L} couplings λ' and \mathcal{B} couplings λ'' are given by

$$\mathcal{L}_{\lambda'} = -\lambda'_{ijk} \left[\tilde{\nu}_L^i \bar{d}_R^k d_L^j + \tilde{d}_L^j \bar{d}_R^k \nu_L^i + (\tilde{d}_R^k)^* (\bar{\nu}_L^i)^c d_L^j - \tilde{e}_L^i \bar{d}_R^k u_L^j - \tilde{u}_L^j \bar{d}_R^k e_L^i - (\tilde{d}_R^k)^* (\bar{e}_L^i)^c u_L^j \right] + h.c., \quad (2)$$

$$\mathcal{L}_{\lambda''} = -\frac{1}{2} \lambda''_{ijk} \left[\tilde{d}_R^k (\bar{u}_R^i)^c d_R^j + \tilde{d}_R^j (\bar{d}_R^k)^c u_R^i + \tilde{u}_R^i (\bar{d}_R^j)^c d_R^k \right] + h.c. \quad (3)$$

With \mathcal{B} couplings, the decay $t \rightarrow ch$ can proceed through the loop diagrams shown in Fig. 1,

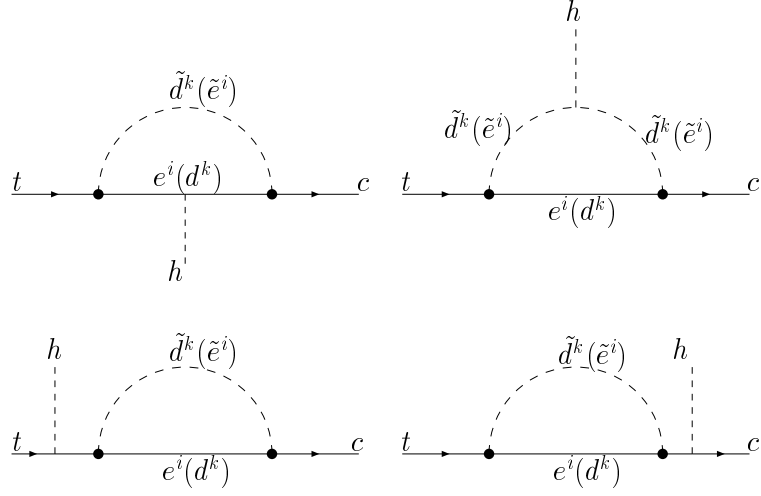


FIG. 2. Feynman diagrams of $t \rightarrow ch$ induced by \mathbb{L} SUSY interactions.

where a down-type quark and a squark are involved in the loops. If the lepton number is violated instead, the decay $t \rightarrow ch$ can occur through the diagrams in Fig. 2, where a down-type quark and a slepton (or a down-type squark and a charged lepton) are in the loops. We note that in the former case the products of \mathcal{B} couplings $\lambda''_{2jk}\lambda''_{3jk}$ are involved while in the latter case the products of \mathbb{L} couplings $\lambda'_{i2k}\lambda'_{i3k}$ are involved.

The induced effective tch vertex is given by, in the cases of \mathcal{B} and \mathbb{L} respectively,

$$V_{\text{tch}}^{\mathcal{B}} = ie \frac{\lambda''_{2jk}\lambda''_{3jk}}{16\pi^2} \xi_c \left(F_L^{\mathcal{B}} P_L + F_R^{\mathcal{B}} P_R \right), \quad (4)$$

$$V_{\text{tch}}^{\mathbb{L}} = ie \frac{\lambda'_{i2k}\lambda'_{i3k}}{16\pi^2} \left(F_L^{\mathbb{L}} P_L + F_R^{\mathbb{L}} P_R \right), \quad (5)$$

where $\xi_c = (N_c - 1)! = 2$ is the color factor and $P_{R,L} = \frac{1}{2}(1 \pm \gamma_5)$. The factors $F_{L,R}^{\mathcal{B}}$ from the contributions of Fig. 1 are given by

$$\begin{aligned} F_L^{\mathcal{B}} &= Y_{dj} m_t m_{dj} (C_0 + 2C_{11}) (-p_t, p_c, m_{dj}, m_{\tilde{d}^k}, m_{dj}) \\ &\quad + Y_{\tilde{d}_R^k} m_t (C_{11} - C_{12}) (-p_t, k, m_{dj}, m_{\tilde{d}^k}, m_{\tilde{d}^k}) \\ &\quad - Y_c \frac{m_t m_c}{m_t^2 - m_c^2} B_1(p_t, m_{dj}, m_{\tilde{d}^k}) - Y_t \frac{m_c^2}{m_c^2 - m_t^2} B_1(p_c, m_{dj}, m_{\tilde{d}^k}), \end{aligned} \quad (6)$$

$$\begin{aligned} F_R^{\mathcal{B}} &= -Y_{dj} m_c m_{dj} (C_0 + 2C_{12}) (-p_t, p_c, m_{dj}, m_{\tilde{d}^k}, m_{dj}) \\ &\quad + Y_{\tilde{d}_R^k} m_c C_{12} (-p_t, k, m_{dj}, m_{\tilde{d}^k}, m_{\tilde{d}^k}) \\ &\quad - Y_c \frac{m_t^2}{m_t^2 - m_c^2} B_1(p_t, m_{dj}, m_{\tilde{d}^k}) - Y_t \frac{m_t m_c}{m_c^2 - m_t^2} B_1(p_c, m_{dj}, m_{\tilde{d}^k}). \end{aligned} \quad (7)$$

The factors $F_{L,R}^{\mathbb{L}} = F_{L,R}^{\mathbb{L}}(\tilde{d}) + F_{L,R}^{\mathbb{L}}(\tilde{e})$, where $F_{L,R}^{\mathbb{L}}(\tilde{d})$ and $F_{L,R}^{\mathbb{L}}(\tilde{e})$ are respectively the contributions of Fig. 2 by exchanging squarks \tilde{d}^k and sleptons \tilde{e}^i , are obtained by the substitutions

$$F_{L,R}^{\mathbb{L}}(\tilde{d}) = F_{R,L}^{\mathcal{B}}|_{dj \rightarrow e^i}, \quad (8)$$

$$F_{L,R}^{\mathbb{L}}(\tilde{e}) = F_{R,L}^{\mathcal{B}}|_{dj \rightarrow d^k, \tilde{d}_R^k \rightarrow \tilde{e}_L^i, m_{dj} \rightarrow -m_{d^k}}. \quad (9)$$

TABLE I. Current 2σ bounds on λ''_{2jk} , λ''_{3jk} , λ'_{i2k} and λ'_{i3k} .

couplings	bounds	sources
$\lambda'_{121}, \lambda'_{122}, \lambda'_{123}$	$0.043 \times (m_{\tilde{d}_R}/100 \text{ GeV})$	charged current universality [15]
λ'_{221}	$0.18 \times (m_{\tilde{d}_R}/100 \text{ GeV})$	ν_μ deep inelastic scattering [15,16]
$\lambda'_{222}, \lambda'_{223}$	$0.21 \times (m_{\tilde{d}_R}/100 \text{ GeV})$	D -meson decay [16]
$\lambda'_{321}, \lambda'_{322}, \lambda'_{323}$	$0.52 \times (m_{\tilde{d}_R}/100 \text{ GeV})$	D_s decay [16],
λ'_{131}	$0.019 \times (m_{\tilde{t}_L}/100 \text{ GeV})$	atomic parity violation [15,17]
λ'_{132}	$0.28 \times (m_{\tilde{t}_L}/100 \text{ GeV})$	asymmetry in e^+e^- collision [15]
λ'_{133}	$0.0014 \times \sqrt{m_{\tilde{b}}}/100 \text{ GeV}$	neutrino mass [18]
λ'_{231}	$0.18 \times (m_{\tilde{b}_L}/100 \text{ GeV})$	ν_μ deep inelastic scattering [15,16]
λ'_{232}	0.56	Z decays [19]
λ'_{233}	$0.15 \times \sqrt{m_{\tilde{b}}}/100 \text{ GeV}$	neutrino mass [18]
$\lambda'_{331}, \lambda'_{332}, \lambda'_{333}$	0.45	Z decays [19]
$\lambda''_{212}, \lambda''_{213}, \lambda''_{223}$	1.23	perturbativity [28]
$\lambda''_{312}, \lambda''_{313}, \lambda''_{323}$	1.0	Z decays [19]

In the above expressions the sum over family indices $i, j, k = 1, 2, 3$ is implied. p_t , p_c and k are the momenta of the top quark, charm quark and Higgs boson, respectively. The functions B_1 and C_{ij} are the conventional 2- and 3-point Feynman integrals [25], and their functional arguments are indicated in the bracket following them. The constants like Y_{di} are the Yukawa couplings of the corresponding particles given by

$$Y_d = \frac{m_d \sin \alpha}{2m_W \sin \theta_W \cos \beta}, \quad (10)$$

$$Y_u = \frac{m_u \cos \alpha}{2m_W \sin \theta_W \sin \beta}, \quad (11)$$

$$Y_{\tilde{d}_R} = -\frac{1}{3}m_Z \tan \theta_W \sin(\alpha + \beta) + \frac{m_d^2}{m_W \sin \theta_W} \frac{\sin \alpha}{\cos \beta}, \quad (12)$$

$$Y_{\tilde{e}_L} = -\frac{m_Z}{\cos \theta_W \sin \theta_W} \left(\frac{1}{2} - \sin^2 \theta_W \right) \sin(\alpha + \beta) + \frac{m_e^2}{m_W \sin \theta_W} \frac{\sin \alpha}{\cos \beta}, \quad (13)$$

where θ_W is the weak mixing angle, α the neutral Higgs boson mixing angle, and $\tan \beta = v_2/v_1$ the ratio of the two vacuum expectation values in the Higgs sector [12].

The ultraviolet divergences are contained in the Feynman integral B_1 . It is easy to check that all the ultraviolet divergences cancel as a result of renormalizability of MSSM.

As a good approximation of neglecting all the fermion masses but m_t , the expression can be substantially simplified. The only contributing diagram in this approximation is that of the Higgs boson coupled to the sfermion. The partial decay width is then given by a simple form

$$\Gamma(t \rightarrow ch) = \frac{K^2(m_t^2 - m_h^2)^2}{32\pi m_t} \left| \frac{(C_{11} - C_{12})(-p_t, k, 0, m_{\tilde{f}}, m_{\tilde{f}})}{16\pi^2} \right|^2, \quad (14)$$

where

$$K = \begin{cases} e \xi_c Y_{\tilde{d}_R^k} \lambda''_{2jk} \lambda''_{3jk} & \text{for } \mathcal{B} \\ e Y_{\tilde{e}_L^i} \lambda'_{i2k} \lambda'_{i3k} & \text{for } \mathcal{L}. \end{cases}$$

This expression is accurate at a level of a few percent.

Now we present the numerical results for $Br(t \rightarrow ch)$ with the full expression. For the SM parameters involved, we take $m_t = 175$ GeV, $m_Z = 91.187$ GeV, $m_W = 80.3$ GeV, $\alpha = 1/128$, $\sin^2 \theta_W = 0.232$, $m_c = 1.7$ GeV, $m_b = 4.7$ GeV, and $m_s = 0.17$ GeV. The results are not sensitive to the light fermion masses. The SUSY parameters involved are the following:

- (1) \mathcal{B} couplings λ''_{2jk} and λ''_{3jk} , and \mathcal{L} couplings λ'_{i2k} and λ'_{i3k} : The current bounds for all \mathcal{R} couplings are summarized in Ref. [21]. In Table 1 we only list those relevant to our study. We note that the current bounds for the \mathcal{B} couplings are generally quite weak.
- (2) Slepton and down-type squark masses: The off-diagonal terms in the sfermion-mass matrices are proportional to the mass of the corresponding fermion, and they are thus relatively small for the down-type squarks and sleptons. Only for large $\tan \beta$, the mass splitting of the sbottoms (staus) becomes sizeable. Nevertheless, the inclusion of the mass splittings for sbottom (stau) does not affect our results appreciably. In our numerical illustration, we thus assume a common mass for all squarks (sleptons).
- (3) Higgs masses and mixing angles of the Higgs sector: At tree level the Higgs sector of the MSSM is determined by two free parameters, *e. g.*, the CP-odd Higgs boson mass m_A and $\tan \beta$. When radiative corrections are included, several other parameters enter through the loops [13]. In phenomenological analysis it is usually to assume a common squark mass scale $m_{\tilde{q}}$ and thus the dominant one-loop corrections can be parameterized by a single parameter ϵ [26],

$$\epsilon = \frac{3g^2}{8\pi^2 m_W^2} m_t^4 \ln \left(1 + \frac{m_{\tilde{q}}^2}{m_t^2} \right). \quad (15)$$

Then the Higgs masses and mixing angle α are obtained by

$$m_{H,h}^2 = \frac{1}{2} \left[m_A^2 + m_Z^2 + \epsilon / \sin^2 \beta \right] \pm \frac{1}{2} \left\{ \left[(m_A^2 - m_Z^2) \cos 2\beta + \epsilon / \sin^2 \beta \right]^2 + (m_A^2 + m_Z^2)^2 \sin^2 2\beta \right\}^{1/2}, \quad (16)$$

$$\tan 2\alpha = \frac{(m_A^2 + m_Z^2) \sin 2\beta}{(m_A^2 - m_Z^2) \cos 2\beta + \epsilon / \sin^2 \beta}. \quad (17)$$

In our calculation we adopt the standard convention [12,26], in which $-\pi/2 \leq \alpha \leq 0$ and $0 \leq \beta \leq \pi/2$ so that as $m_A \rightarrow \infty$, we have $\alpha \rightarrow \beta - \pi/2$ and thus the h couplings approach to those of the SM Higgs boson while other Higgs bosons (A, H, H^\pm) decouple.

The current constraints on the parameter space of the Higgs sector were given by the LEP experiments [27]. The region of small $\tan \beta$ and low Higgs mass m_A has

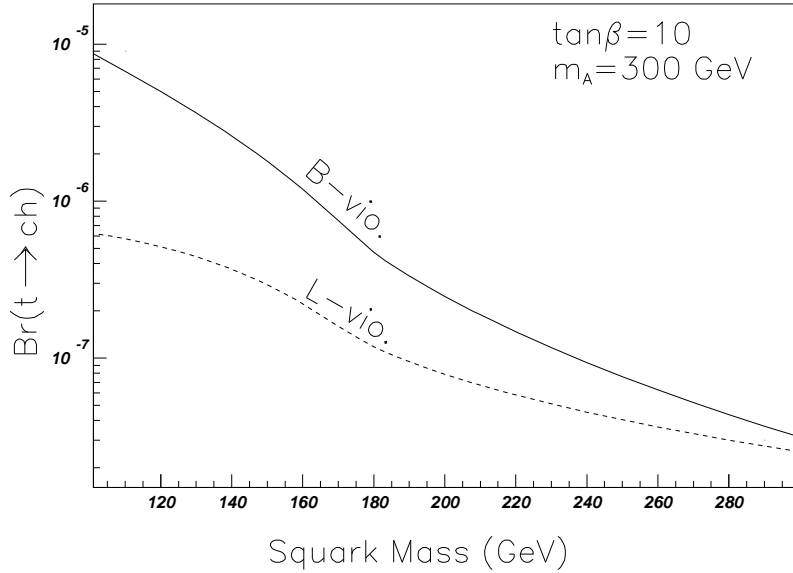


FIG. 3. $Br(t \rightarrow ch)$ versus squark mass for $m_A = 300$ GeV and $\tan\beta = 10$. The solid and dashed curves correspond to B-violation and L-violation, respectively.

been excluded, depending on the extent of top-squark mixing. In any case the robust bounds $m_h > 83.4$ GeV and $m_A > 83.8$ GeV have been obtained at 95% confidence level for $\tan\beta$ greater than 0.8.

The decay rate increases quadratically with the relevant λ' or λ'' couplings and decreases with the increase of the sparticle mass. Using the upper limits of the relevant \mathcal{R} couplings in Table 1, we present the maximum values of the branching fraction in Figs. 3 – 5.

Fig. 3 shows the dependence of $Br(t \rightarrow ch)$ on squark mass for $m_A = 300$ GeV and $\tan\beta = 10$. The solid and dashed curves correspond to B-violation and L-violation, respectively. With the parameter values in this figure, the value of m_h is within the range $93 \sim 106$ GeV corresponding to the squark mass range $100 \sim 300$ GeV. The results manifest the decoupling property of the MSSM, *i. e.*, the contributions drop with the increase of the squark mass. We found that when squarks become sufficiently heavy ($\gtrsim 500$ GeV), they decouple quickly and the branching fraction goes down like $1/m_q^4$ asymptotically. From Fig. 3 one sees that in the \mathcal{B} case, for a squark as light as 100 GeV, the branching fraction can reach the level of 10^{-5} . In the \mathcal{L} case, however, the contributions are below the level of 10^{-6} because of the current stringent bounds on the relevant couplings in Table 1. Thus in the following we do not present the \mathcal{L} contributions any more.

Fig. 4 shows the \mathcal{B} contributions to $Br(t \rightarrow ch)$ as a function of $\tan\beta$ for $m_A = 300$ GeV. The solid and dashed curves correspond to squark mass of 100 GeV and 200 GeV, respectively. The hatched region is excluded by LEP experiments [27]. Note that in the case of maximal top-squark mixing, the very small $\tan\beta$ region ($\lesssim 0.2$) is still allowed by LEP experiments and in this region the branching fraction can also be significant. Regarding m_h , in the small $\tan\beta$ range $0.01 \lesssim \tan\beta \lesssim 0.2$, we have $91 \lesssim m_h \lesssim 92$ GeV ($91 \lesssim m_h \lesssim 93$) for squark mass of 100 GeV (200 GeV). In the large $\tan\beta$ range $4 \lesssim \tan\beta \lesssim 100$, we

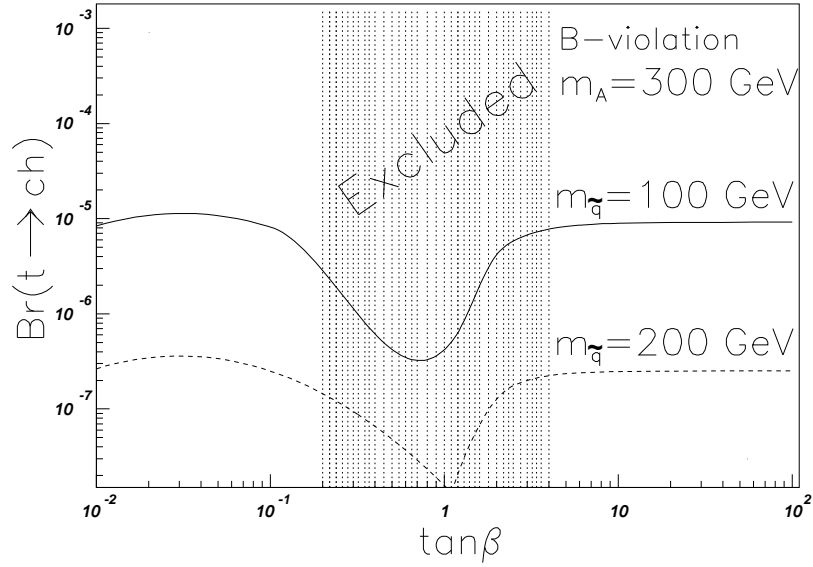


FIG. 4. \mathcal{B} contributions to $Br(t \rightarrow ch)$ versus $\tan \beta$ for $m_A = 300$ GeV. The solid and dashed curves correspond to squark mass of 100 GeV and 200 GeV, respectively. The hatched region is excluded by LEP experiments.

have $84 \lesssim m_h \lesssim 94$ GeV ($91 \lesssim m_h \lesssim 101$ GeV) for squark mass of 100 GeV (200 GeV), correspondingly.

Fig. 5 is the \mathcal{B} contribution to $Br(t \rightarrow ch)$ versus m_A for $\tan \beta = 10$. The solid and dashed curves correspond to squark mass of 100 GeV and 200 GeV, respectively. The range of m_h is $83 \lesssim m_h \lesssim 93$ GeV and $85 \lesssim m_h \lesssim 100$ GeV for squark mass of 100 GeV, 200 GeV, respectively. This figure shows that the dependence of the branching fraction on m_A is quite mild.

A couple of remarks are in order regarding the above results:

- (1) The branching fraction is very sensitive to the relevant \mathcal{R} couplings, *e. g.*, in the \mathcal{B} case, is proportional to $(\lambda''_{2jk}\lambda''_{3jk})^2$. The relaxation of some bounds in Table 1 will result in large enhancement of the branching fraction. For example, so far no experimental bound for λ''_{2jk} has been reported and we used the theoretical perturbativity bound ($\simeq 1.23$) [28]. Such a bound is derived from the assumption that the gauge groups unify at $M_U = 2 \times 10^{16}$ GeV and the Yukawa couplings Y_t, Y_b and Y_τ remain in the perturbative domain in the whole range up to M_U , *i. e.*, $Y_i(\mu) < 1$ for $\mu < 2 \times 10^{16}$ GeV. But there is no a priori reason to take this theoretical assumption. So the bound on λ''_{2jk} could be even weaker.
- (2) The decay rate is proportional to a product of two \mathcal{R} couplings, and thus the non-zero decay requires the coexistence of at least two couplings. Although we do not expect any stringent constraints on the coupling products involving the third-generation fermions, it would be interesting to seek for other possible FCNC processes at low energies which may be induced by such coexistence of more than one \mathcal{R} coupling.

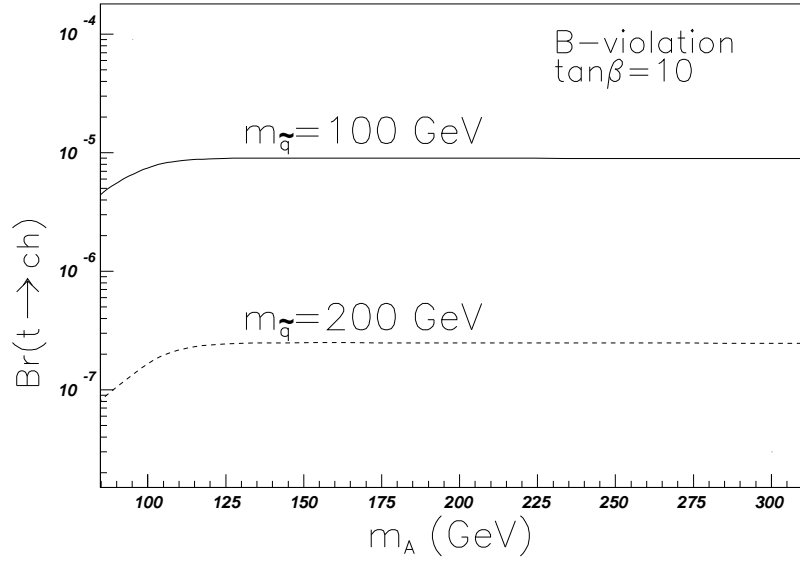


FIG. 5. \mathcal{B} contributions to $Br(t \rightarrow ch)$ versus m_A for $\tan \beta = 10$. The solid and dashed curves correspond to squark mass of 100 GeV and 200 GeV, respectively.

- (3) We assumed a common mass for all squarks or sleptons for illustration of the numerical results. Our results are not very sensitive to the mass differences of the sfermions because the involved \mathcal{R} couplings, λ'' or λ' , are not unitary in general so that the delicate cancellations (like the GIM mechanism) do not generally occur.
- (4) Although in our analysis we varied $\tan \beta$ freely and only considered the direct LEP constraints, we note that the extreme case of a very small or very large $\tan \beta$ value may be severely disfavored by some theoretical arguments and/or by some indirect experimental constraints like those from $b \rightarrow s\gamma$ [29]. Since the branching ratio of $t \rightarrow ch$ could be significantly enhanced also in the intermediate range of $\tan \beta$ (say $4 \lesssim \tan \beta \lesssim 30$), our conclusion will not be affected by such indirect experimental constraints.
- (5) Finally, we point out that we only presented the analysis for the most interesting decay channel $t \rightarrow ch$ because kinematically it is always allowed. Other decay channels, $t \rightarrow cH$ or $t \rightarrow cA$, are quite likely to be suppressed by phase space or even kinematically forbidden. If the channel $t \rightarrow cA$ is kinematically allowed, the branching ratio $Br(t \rightarrow cA)$ can also be significant for a very large $\tan \beta$ ($\gtrsim 40$) because its coupling to bottom quark is proportional to $\frac{m_b}{m_W} \tan \beta$. Since its coupling to top quark is proportional to $\frac{m_t}{m_W} \cot \beta$, $Br(t \rightarrow cA)$ could be enhanced significantly for a small $\tan \beta$. However, for such a light Higgs boson A ($m_A \lesssim 175$ GeV), the small $\tan \beta$ region has been almost completely excluded by LEP experiments [27].

In summary, we found that the FCNC decay $t \rightarrow ch$ can be significantly enhanced relative to that of the SM in SUSY theories with R-parity violation. The branching fraction

depends quadratically upon the products of \mathcal{R} couplings, and scales with the heavy mass of the sfermions in the loops as $m_{\tilde{f}}^{-4}$, and it can be at least as large as that in the MSSM. In the optimistic scenario that the involved couplings take their current upper bounds and $m_{\tilde{f}} \approx 100$ GeV, the branching fraction can be as large as 10^{-5} , which is potentially accessible at the LHC with a high luminosity.

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